

MZ-TH/99-45

**Λ_b and Λ_c baryon decays at finite values of
heavy quark masses**

M. A. Ivanov

*Bogoliubov Laboratory of Theoretical Physics, Joint Institute for Nuclear Research,
141980 Dubna, Russia*

J. G. Körner

Johannes Gutenberg-Universität, Institut für Physik, D-55099 Mainz, Germany

V. E. Lyubovitskij

*Bogoliubov Laboratory of Theoretical Physics, Joint Institute for Nuclear Research,
141980 Dubna, Russia and Department of Physics, Tomsk State University,
634050 Tomsk, Russia*

A. G. Rusetsky

*Institute for Theoretical Physics, University of Bern, Sidlerstrasse 5, CH-3012,
Bern, Switzerland and Bogoliubov Laboratory of Theoretical Physics, Joint
Institute for Nuclear Research, 141980 Dubna, Russia and HEPI, Tbilisi State
University, 380086 Tbilisi, Georgia*

Abstract

Semileptonic decays of Λ_b and Λ_c baryons are studied within the Relativistic Three-Quark Model using finite heavy quark mass values. Employing the same parameters as have been used previously for the description of exclusive decays of heavy baryons in the heavy quark limit we calculate the six form factors of the process and the corresponding decay rates. Our calculation shows that the “finite mass” corrections are important in heavy-to-light transitions and are not negligible in heavy-to-heavy transitions.

Pacs: 11.15.Me, 11.30.Hv, 12.38.Lg, 12.39.Ki, 14.20.Lq, 14.20.Mr

Keywords: relativistic quark model, charm and bottom baryons, semileptonic decays, finite quark masses, heavy quark limit

During the last decade heavy baryon transitions (semileptonic, nonleptonic, strong and electromagnetic) have been studied in detail within the Heavy Quark Effective Theory employing QCD sum rule methods or nonrelativistic and relativistic quark models, etc. (see, for example, the reviews in [1,2] and the papers [3–17]). The mass spectrum of heavy baryons as well as their exclusive and inclusive decays have been described successfully in these approaches incorporating the ideas of QCD.

In the papers [12–14,18–21] we proposed and developed a QCD motivated Relativistic Three-Quark Model (RTQM), which can be viewed as an effective quantum field approach based on an interaction Lagrangian of light and heavy baryons interacting with their constituent quarks. The coupling strength of the baryons interacting with the three constituent quarks are determined by the compositeness condition $Z_H = 0$ [19,22] where Z_H is the wave function renormalization constant of the hadron. The compositeness condition enables one to unambiguously and consistently relate theories with quark and hadron degrees of freedom to the effective Lagrangian approaches formulated in terms of hadron variables only (as, for example, Chiral Perturbation Theory [23] and its covariant extension to the baryon sector [24]). Our strategy is as follows. We start with an effective interaction Lagrangian written down in terms of quark and hadron variables. Then, by using Feynman rules, the S -matrix elements describing hadron-hadron interactions are given in terms of a set of quark diagrams. The compositeness condition enables one to avoid a double counting of quark and hadron degrees of freedom. The RTQM model contains only a few model parameters: the masses of the light and heavy quarks, and certain scale parameters that define the size of the distribution of the constituent quarks inside the hadron. The RTQM approach has been previously used to compute the exclusive semileptonic, nonleptonic, strong and electromagnetic decays of charm and bottom baryons [12–14] in the heavy quark limit $m_Q \rightarrow \infty$ always employing the same set of model parameters.

The main objective of this letter is to extend our approach to the study of heavy baryon transitions at finite values of the heavy quark mass without using an explicit $1/m_Q$ expansion. In the following we confine ourself to the dominant exclusive semileptonic transitions of Λ_b and Λ_c baryons: $\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e$ and $\Lambda_c \rightarrow \Lambda_s e^+ \nu_e$. We compare our results to a similar finite mass calculation done in [17], where the QCD sum rule approach was used. In agreement with [17] we find that finite quark mass corrections to the form factors and the rates of semileptonic transitions are important for heavy-to-light $c \rightarrow s$ transitions and not negligible for heavy-to-heavy $b \rightarrow c$ transitions.

We proceed as follows. First we briefly explain the basic ideas of the Relativistic Three-Quark Model (RTQM) and describe our calculational techniques for the finite heavy quark case. We then compute form factors and rates for the decays $\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e$ and $\Lambda_c \rightarrow \Lambda_s e^+ \nu_e$. For reasons of comparison numeri-

cal results are provided both for the finite quark mass case and the infinite quark mass case. We finally compare our results with the results of recent calculations within the QCD sum rule method [17].

We start with a brief description of our approach, the Relativistic Three-Quark Model (RTQM). A detailed description of the RTQM can be found in Refs. [12,14,20,21]. In the RTQM approach baryons are described as bound states of constituent quarks. The coupling of the baryons with their constituents is defined through an effective relativistic interaction Lagrangian which contains the usual three-quark currents with the quantum numbers of the heavy baryons [12]. For example, the Lagrangian of the Λ -type baryons (Λ_b , Λ_c and Λ_s) coupling to three quarks is taken as

$$\begin{aligned} \mathcal{L}_{\Lambda_q}^{\text{int}}(x) &= g_{B_q} \bar{\Lambda}_q(x) \Gamma_1 q^a(x) \int d\xi_1 \int d\xi_2 F_{B_q}(\xi_1^2 + \xi_2^2) \\ &\quad \times u^b(x + 3\xi_1 - \sqrt{3}\xi_2) C \Gamma_2 d^c(x + 3\xi_1 + \sqrt{3}\xi_2) \varepsilon^{abc} + \text{h.c.} \\ F_{B_q}(\xi_1^2 + \xi_2^2) &= \int \frac{d^4 k_1}{(2\pi)^4} \int \frac{d^4 k_2}{(2\pi)^4} e^{ik_1 \xi_1 + ik_2 \xi_2} \tilde{F}_{B_q}(k_1^2 + k_2^2) \end{aligned} \quad (1)$$

where the $q = b, c, s$ stands for the quark spinor and Γ_i are spinor matrices which define the quantum numbers of the relevant three-quark currents. $C = \gamma^0 \gamma^2$ is the charge conjugation matrix and g_{B_q} is the coupling constant which is determined by the compositeness condition [14,19,22]. The compositeness condition implies that the renormalization constant of the hadron wave function is set equal to zero: $Z_{B_q} = 1 - g_{B_q}^2 \Sigma'_{B_q}(M_{B_q}) = 0$ where Σ'_{B_q} is the derivative of the baryon mass operator and M_{B_q} is the baryon mass.

It is well known that the form of the three-quark current for baryons is not unique. In fact, one can write down different interpolating currents for a given baryon (even in the absence of derivative interactions) that have the correct quantum numbers of the given baryon. For finite mass Λ -type baryons there are three possibilities for the choice of the three-quark current without derivatives. In papers [18,25] we showed that baryon observables are only weakly dependent on the choice of the three-quark currents. In this paper we consider the simplest set of baryon currents: the pseudoscalar currents for the heavy baryons Λ_b and Λ_c ($\Gamma_1 \otimes C \Gamma_2 = I \otimes C \gamma_5$) [12] and a $SU(3)$ symmetric (tensor) current for the Λ_s hyperon ($\Gamma_1 \otimes C \Gamma_2 = I \otimes C \gamma_5 + \gamma_5 \otimes C$) [18].

In Eq. (1) we have introduced a baryon-three-quark vertex form factor given by $\tilde{F}_{B_q}(k_1^2 + k_2^2)$. For simplicity, we factorize out the $q = b, c$ or s quark in the Λ -type baryons by placing them at the center of the baryon. This means that we take the masses of the b , c and s quarks to be much larger than the masses of u and d quarks. Any choice of vertex function F_{B_q} is appropriate as long as it falls off sufficiently fast in the ultraviolet region to render the Feynman diagrams

ultraviolet finite. In principle, their functional forms would be calculable from the solutions of the Bethe-Salpeter equations for the baryon bound states [15] which is, however, an untractable problem at present. In our previous analysis [20] we found that, using various forms for the vertex function, the hadron observables are insensitive to the exact details of the functional form of the hadron-quark vertex form factor. We will use this observation as a guiding principle and choose simple Gaussian forms for the vertices F_{B_q} . Their Fourier transform reads [12–14,21]

$$\tilde{F}_{B_q}(k_1^2 + k_2^2) = \exp\left(\frac{k_1^2 + k_2^2}{\Lambda_{B_q}^2}\right) \quad (2)$$

where Λ_{B_q} is a scale parameter defining the distribution of the u and d quarks in Λ -type baryon. For the light and heavy quark propagators with constituent masses M we shall use the standard form of the free fermion propagator

$$S_M(k) = \frac{1}{M - \not{k}} \quad (3)$$

where $M = m$ for the u or d quarks, $M = m_s$ for the strange quark, $M = m_c$ for the charm quark and $M = m_b$ for the bottom quark.

Next we specify our model parameters. In order to be able to compare the two calculations with finite quark masses on the one hand and infinite quark masses [14] on the other hand we use the same set of model parameters in both calculations: i) $\Lambda_B = \Lambda_{B_u} = \Lambda_{B_d} = \Lambda_{B_s} = 1$ GeV is the common scale parameter defining the distribution of quarks in light baryons [21]; ii) in the heavy flavour sector the scale parameters Λ_{B_c} and Λ_{B_b} are chosen to be the same $\Lambda_{B_Q} = \Lambda_{B_c} = \Lambda_{B_b} = 1.8$ GeV in order to provide the correct normalization of the baryonic Isgur-Wise function in the heavy quark limit [14]; iii) the values of the constituent quark masses are fixed from the analysis of magnetic moments and charge radii of light baryons [14,21]: $m = 420$ MeV and $m_s = 570$ MeV. Therefore, we have only two free parameters: m_c and m_b , the masses of the charm and bottom quark. Their values are fixed according to [14]

$$\bar{\Lambda} = M_{\Lambda_b} - m_b = M_{\Lambda_c} - m_c = 600 \text{ MeV} \quad (4)$$

In this way one can meaningfully compare the results of the two calculations for finite and infinite quark mass values. Using Eq. (4) with the experimental values for the Λ_b and Λ_c baryon masses $m_{\Lambda_b} = 5.64$ GeV and $m_{\Lambda_c} = 2.285$ GeV [26] we find $m_b = 5.04$ GeV and $m_c = 1.685$ GeV. The results obtained in this paper will be compared to the results of [14] done in the heavy quark limit ($m_Q \rightarrow \infty$). In both cases the difference between the masses of the heavy

baryons and the masses of the heavy quarks is fixed at $\bar{\Lambda} = 600$ MeV. The two calculations differ only in the choice of the heavy quark propagators: in the finite mass scheme we use the usual free propagator (without any $1/m_Q$ expansion) and in the infinite mass scheme we use the usual leading HQET propagator [12,14]

$$S_v(k, \bar{\Lambda}) = -\frac{(1+\not{v})}{2(v \cdot k + \bar{\Lambda})} \quad (5)$$

where the four-velocity of the heavy quark is denoted by v as usual.

The semileptonic decays of the Λ_b and Λ_c baryons are described by the triangle two-loop diagram shown Fig.1. Correspondingly one has

$$M_\mu(\Lambda_{q_1} \rightarrow \Lambda_{q_2}) = \bar{u}(p_2)\Lambda_\mu(p_1, p_2)u(p_1) = \frac{g_{B_{q_1}}}{8\pi^2} \frac{g_{B_{q_2}}}{8\pi^2} \bar{u}(p_2)\Gamma_\mu(p_1, p_2)u(p_1) \quad (6)$$

for the matrix element of the transition. The weak vertex function $\Gamma_\mu(p_1, p_2)$ takes the form

$$\begin{aligned} \Gamma_\mu(p_1, p_2) = & \int \frac{d^4 k_1}{\pi^2 i} \int \frac{d^4 k_2}{4\pi^2 i} \tilde{F}_{B_{q_1}}(12[k_1^2 + k_1 k_2 + k_2^2]) \tilde{F}_{B_{q_2}}(12[k_1^2 + k_1 k_2 + k_2^2]) \\ & \times \Gamma_1^f S_{q_2}(k_1 + p_2) O_\mu S_{q_1}(k_1 + p_1) \Gamma_1^i \text{Tr}[\Gamma_2^f S_{q_4}(k_2) \Gamma_2^i S_{q_3}(k_1 + k_2)] \end{aligned} \quad (7)$$

where $\Gamma_{1(2)}^i$ and $\Gamma_{1(2)}^f$ are the Dirac matrices of the initial and the final baryons, respectively, and $O_\mu = \gamma_\mu(1 + \gamma_5)$. The integral (7) is calculated in the Euclidean region both for internal and external momenta. The final results are obtained by analytic continuation of the external momenta to the physical region after the internal momenta have been integrated out. In order to keep the calculation as general as possible we shall retain a general form for the vertex function \tilde{F}_{B_q} in our analytical integrations. Only at the very end of the calculation when we do the numerical evaluation the Gaussian form (2) will be substituted for the vertex function. As mentioned before the difference to the earlier calculation in Ref. [14] lies in the use of the full quark propagator, whereas in [14] we have used the leading HQET propagator. The integration techniques used in [14] can be easily extended to the case of finite heavy quark masses.

As an illustration of our calculational procedure we evaluate integral (7) for equal values of the heavy and light scale parameters $\Lambda_{B_{q_1}} = \Lambda_{B_{q_2}} = \Lambda_{B_Q}$. First of all, we express all dimensional parameters entering the two-loop integral (7) in units of Λ_{B_Q} . We then have

$$\begin{aligned}
\Gamma_\mu(p_1, p_2) = & \int_0^\infty ds \tilde{F}_B^L(s) \int_0^\infty d^4\alpha \ e^{-m_3^2\alpha_3 - m_4^2\alpha_4 - (m_1^2 - p_1^2)\alpha_1 - (m_2^2 - p_2^2)\alpha_2} \\
& \times \Gamma_1^f \left(m_2 + \not{p}_2 + \frac{\not{\partial}_1}{2} \right) O_\mu \left(m_1 + \not{p}_1 + \frac{\not{\partial}_1}{2} \right) \Gamma_1^i \\
& \times \text{Tr} \left[\Gamma_2^f \left(m_4 + \frac{\not{\partial}_2}{2} \right) \Gamma_2^i \left(m_3 + \frac{\not{\partial}_1}{2} + \frac{\not{\partial}_2}{2} \right) \right] \int \frac{d^4 k_1}{\pi^2 i} \int \frac{d^4 k_2}{4\pi^2 i} e^{k \mathcal{A} k + 2k B}
\end{aligned} \tag{8}$$

where $\not{\partial}_i = \gamma_\mu \cdot \partial / \partial B_i^\mu$, $\tilde{F}_B^L(s)$ is the Laplace transform of $\tilde{F}_{B_Q}^2(12[k_1^2 + k_1 k_2 + k_2^2])$ and the matrices \mathcal{A} and B are defined by

$$\begin{aligned}
\mathcal{A}_{ij} = & \begin{pmatrix} 12s + \alpha_1 + \alpha_2 + \alpha_3 & 6s + \alpha_3 \\ 6s + \alpha_3 & 12s + \alpha_3 + \alpha_4 \end{pmatrix} \\
B_i = & \begin{pmatrix} \alpha_1 p_1 + \alpha_2 p_2 \\ 0 \end{pmatrix}
\end{aligned}$$

The integration over k_1 , k_2 and the variable s results in

$$\begin{aligned}
\Gamma_\mu(p_1, p_2) = & \int_0^\infty d^4\alpha \left\{ \frac{\tilde{F}_{B_Q}^2(-12w)}{|A|^2} \Gamma_1^f (m_2 + \not{p}_2 - \not{B}_1 A_{11}^{-1}) O_\mu \right. \\
& \times (m_1 + \not{p}_1 - \not{B}_1 A_{11}^{-1}) \Gamma_1^i \text{Tr}[\Gamma_2^f (m_4 - \not{B}_1 A_{12}^{-1}) \Gamma_2^i (m_3 - \not{B}_1 (A_{11}^{-1} + A_{12}^{-1}))] \\
& - \int_0^\infty dt \frac{\tilde{F}_{B_Q}^2(-12[w+t])}{2|A|^2} \left[(A_{12}^{-1} + A_{22}^{-1}) \text{Tr}[\Gamma_2^f \gamma^\beta \Gamma_2^i \gamma_\beta] \Gamma_1^f (m_2 + \not{p}_2 - \not{B}_1 A_{11}^{-1}) \right. \\
& \times O_\mu (m_1 + \not{p}_1 - \not{B}_1 A_{11}^{-1}) \Gamma_1^i + \Gamma_1^f [\gamma^\alpha O_\mu (m_1 + \not{p}_1 - \not{B}_1 A_{11}^{-1}) \\
& + (m_2 + \not{p}_2 - \not{B}_1 A_{11}^{-1}) O_\mu \gamma^\alpha] \Gamma_1^i \left(A_{12}^{-1} \text{Tr}[\Gamma_2^f \gamma_\alpha \Gamma_2^i (m_3 - \not{B}_1 (A_{11}^{-1} + A_{12}^{-1}))] \right. \\
& + (A_{11}^{-1} + A_{12}^{-1}) \text{Tr}[\Gamma_2^f (m_4 - \not{B}_1 A_{12}^{-1}) \Gamma_2^i \gamma_\alpha] \left. \right) \\
& \left. + \Gamma_1^f \gamma^\alpha O_\mu \gamma_\alpha \Gamma_1^i A_{11}^{-1} \text{Tr}[\Gamma_2^f (m_4 - \not{B}_1 A_{12}^{-1}) \Gamma_2^i (m_3 - \not{B}_1 (A_{11}^{-1} + A_{12}^{-1}))] \right] \\
& + \int_0^\infty dt t \frac{\tilde{F}_{B_Q}^2(-12[w+t])}{4|A|^2} \left[\Gamma_1^f \gamma^\alpha O_\mu \gamma_\alpha \Gamma_1^i A_{11}^{-1} (A_{12}^{-1} + A_{22}^{-1}) \text{Tr}[\Gamma_2^f \gamma^\beta \Gamma_2^i \gamma_\beta] \right. \\
& \left. + \Gamma_1^f \gamma^\alpha O_\mu \gamma^\beta \Gamma_1^i A_{12}^{-1} (A_{11}^{-1} + A_{12}^{-1}) \text{Tr}[\Gamma_2^f \gamma_\alpha \Gamma_2^i \gamma_\beta + \Gamma_2^f \gamma_\beta \Gamma_2^i \gamma_\alpha] \right] \}
\end{aligned} \tag{9}$$

where M_1 and M_2 are the masses of initial and final baryons, respectively, and

$$w = \sum_{i=1}^4 m_i^2 \alpha_i - M_1^2 \alpha_1 - M_2^2 \alpha_2 + A_{11}^{-1} (p_1 \alpha_1 + p_2 \alpha_2)^2,$$

$$A_{ij}^{-1} = \frac{1}{|A|} \begin{pmatrix} 1 + \alpha_3 + \alpha_4 & -(1/2 + \alpha_3) \\ -(1/2 + \alpha_3) & 1 + \alpha_1 + \alpha_2 + \alpha_3 \end{pmatrix}$$

The vertex function $\Lambda_\mu(p_1, p_2)$ is as usual decomposed into a set of six relativistic form factors which are functions of the momentum transfer squared $t = q^2$. We shall present our results in terms of two alternative sets of heavy baryon weak form factors [1,3,6]. The two sets of form factors are defined by the covariant expansions

$$\Lambda_\mu(p_1, p_2) = \gamma_\mu (F_1^V + F_1^A \gamma_5) + i\sigma_{\mu\nu} q^\nu (F_2^V + F_2^A \gamma_5) + q^\mu (F_3^V + F_3^A \gamma_5) \quad (10)$$

and

$$\Lambda_\mu(p_1, p_2) = \gamma_\mu (G_1^V + G_1^A \gamma_5) + v_\mu (G_2^V + G_2^A \gamma_5) + v'_\mu (G_3^V + G_3^A \gamma_5) \quad (11)$$

where $v_\mu = p_\mu/M_1$ and $v'_\mu = p'_\mu/M_2$ are the four-velocities of the initial and final baryon, respectively. The relation between the two sets of heavy baryon form factors can be easily worked out [1,6]. In the heavy quark limit [3] one has

$$F_1^V = F_1^A = G_1^V = G_1^A, \quad F_2^{V(A)} = F_3^{V(A)} = G_2^{V(A)} = G_3^{V(A)} = 0 \quad (12)$$

for the heavy-to-heavy transition $\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e$. For heavy-to-light transitions as in $\Lambda_c \rightarrow \Lambda_s e^+ \nu_e$ transition one has

$$G_2^V = G_2^A = G_1^A - G_1^V, \quad G_3^V = G_3^A = 0, \quad F_1^V = F_1^A, \quad F_2^{V(A)} = F_3^{V(A)} \quad (13)$$

in the heavy (c -quark) limit [4,7]. We shall not write down rate and asymmetry formulae in terms of these form factors since these have been worked out in great detail in Refs. [1,6,8,12].

We now present our numerical results for the exclusive semileptonic decays $\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e$ and $\Lambda_c \rightarrow \Lambda e^+ \nu_e$ for the finite and infinite mass cases. In Table I we present our results for the two sets of relativistic form factors (Set I and Set II)¹ for two values of q^2 , namely $q^2 = q_{max}^2 = (M_1 - M_2)^2$ and $q^2 = 0$.

¹ We use the notation:

In parenthesis we give the values of the form factors in the heavy quark limit. Table I shows that in the case of the $b \rightarrow c$ transition the “finite mass” corrections can amount to $\sim 10\%$ for the Set I form factors and up to $\sim 25\%$ for the Set II form factors. In the case of $c \rightarrow s$ transitions the “finite mass” corrections are significantly larger (see TABLE II). Our estimate of the “finite mass” corrections for $\Lambda_b \rightarrow \Lambda_c$ transition at $q^2 = q_{max}^2$ agrees with the QCD sum rules calculations: $F_1^V = F_1^A = 1.03 \pm 0.06$ and $F_2^{V(A)} = F_3^{V(A)} = 0$ [17]. In TABLE III we show our results for the rates and mean asymmetries of Λ_b and Λ_c decays again for the finite and infinite mass cases. For the heavy-to-heavy transition $\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e$ the finite mass corrections are $\sim 10\%$ in the decay rate and $\sim 7\%$ in the asymmetry parameter. In the case of $\Lambda_c \rightarrow \Lambda_s$ transition one can see that the finite mass corrections are larger than in the $\Lambda_b \rightarrow \Lambda_c$ case. They amount to $\sim 50\%$ for the decay rate and $\sim 10\%$ for the asymmetry parameter. Our prediction for the $\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e$ rate is in agreement with the experimental upper limit given by $\Gamma(\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e) = (6.67 \pm 2.73) \times 10^{10} \text{ sec}^{-1}$ [26]. Our prediction for the $\Lambda_c \rightarrow \Lambda_s e^+ \nu_e$ rate agrees with the corresponding experimental value measured by the CLEO Collaboration: $\Gamma(\Lambda_c \rightarrow \Lambda_s e^+ \nu_e) = (9.54 \pm 2.28) \times 10^{10} \text{ sec}^{-1}$ [26].

In conclusion, we have found that the finite mass corrections significantly contribute to the rate and the asymmetry parameter in the heavy-to-light transition $\Lambda_c \rightarrow \Lambda_s e^+ \nu_e$ whereas the finite mass corrections are smaller for the heavy-to-heavy transition $\Lambda_b \rightarrow \Lambda_c e^- \bar{\nu}_e$. It appears that the leading term in the heavy quark expansion gives a reasonably accurate description of the heavy-to-heavy $b \rightarrow c$ transitions. Contrary to this the finite mass corrections are quite important for the heavy-to-light $c \rightarrow s$ transitions.

M.A.I. and V.E.L. thank Mainz University for hospitality where this work was completed. The visit of M.A.I. to Mainz University was supported by the DFG (Germany) and the visit of V.E.L. to Mainz University was supported by the Graduiertenkolleg ”Eichtheorien”, Mainz. This work was supported in part by the Heisenberg-Landau Program and by the BMBF (Germany) under contract 06MZ865. J.G.K. acknowledges partial support by the BMBF (Germany) under contract 06MZ865. A.G.R. acknowledges partial support of the Swiss National Science Foundation, and TMR, BBW-Contract No. 97.0131 and EC-Contract No. ERBFMRX-CT980169 (EURODAΦNE).

Set I - the set of form factors $(F_1^V, F_2^V, F_3^V, F_1^A, F_2^A, F_3^A)$,
Set II - the set of form factors $(G_1^V, G_2^V, G_3^V, G_1^A, G_2^A, G_3^A)$.

References

- [1] J.G. Körner, M. Krämer, and D. Pirjol, *Prog. in Part. Nucl. Phys.* **33** (1994) 787.
- [2] M. Neubert, *Phys. Rep.* **245** (1994) 259.
- [3] N. Isgur and M. Wise, *Nucl. Phys.* **B348** (1991) 276.
- [4] F. Hussain, J.G. Körner, M. Krämer, and G. Thompson, *Z. Phys.* **C51** (1991) 321.
- [5] T. Mannel, W. Roberts, and Z. Ryzak, *Nucl. Phys.* **B355** (1991) 38.
- [6] G.V. Efimov, M.A. Ivanov, and V.E. Lyubovitskij, *Z. Phys. C* **52** (1991) 149; G.V. Efimov, M.A. Ivanov, N.B. Kulimanova, and V.E. Lyubovitskij, *Z. Phys. C* **54** (1992) 349.
- [7] F. Hussain et. al., *Nucl. Phys.* **B370** (1992) 259.
- [8] J.G. Körner and M. Krämer, *Phys. Lett.* **B275** (1992) 495; *Z. Phys.* **C55** (1992) 659.
- [9] T.M. Yan et. al., *Phys. Rev.* **D46**, 1148 (1992).
- [10] H.-Y. Cheng et al. *Phys. Rev.* **D47**, 1030 (1993).
- [11] M.-Q. Huang, Y.-B. Dai, and C.-S. Huang, *Phys. Rev.* **D52**, 3986 (1995).
- [12] M.A. Ivanov, V.E. Lyubovitskij, J.G. Körner and P. Kroll, *Phys. Rev.* **D56** (1997) 348.
- [13] M.A. Ivanov, J.G. Körner, V.E. Lyubovitskij, and A.G. Rusetsky, *Phys. Rev.* **D57** (1998) 5638; *Mod. Phys. Lett.* **A13** (1998) 181.
- [14] M.A. Ivanov, J.G. Körner, V.E. Lyubovitskij, and A.G. Rusetsky, *Phys. Lett.* **B442** (1998) 435; M.A. Ivanov, J.G. Körner, and V.E. Lyubovitskij, *Phys. Lett.* **B448** (1999) 143; M.A. Ivanov, J.G. Körner, V.E. Lyubovitskij, and A.G. Rusetsky, *Phys. Rev.* **D60** (1999) 094002.
- [15] M.A. Ivanov, J.G. Körner, V.E. Lyubovitskij, and A.G. Rusetsky, *Phys. Rev.* **D59** (1999) 074016.
- [16] S. Tawfiq, P.J. O'Donnell, and J.G. Körner, *Phys. Rev.* **D58** (1998) 054010; S. Tawfiq and P.J. O'Donnell, *Phys. Rev.* **D60** (1999) 014013.
- [17] H.G. Dosch, E. Ferreira, M. Nielsen, and R. Rosenfeld, *Phys. Lett.* **B431** (1998) 173; R.S.M. de Carvalho, et al., Preprint hep-ph/9903326.
- [18] G.V. Efimov, M.A. Ivanov, and V.E. Lyubovitskij, *Few-Body Syst.* **6** (1989) 17.
- [19] G.V. Efimov and M.A. Ivanov, "The Quark Confinement Model", IOP, 1993.

- [20] I.V. Anikin, M.A. Ivanov, N.B. Kulimanova, and V.E. Lyubovitskij, *Z. Phys. C* **65** (1995) 681.
- [21] M.A. Ivanov, M.P. Locher, and V.E. Lyubovitskij, *Few-Body Syst.* **21** (1996) 131.
- [22] A. Salam, *Nuovo Cim.* **25** (1962) 224; S. Weinberg, *Phys. Rev.* **130** (1963) 776; K. Hayashi et al., *Fort. Phys.* **15** (1967) 625.
- [23] S. Weinberg, *Physica* **96A** (1979) 327; J. Gasser and H. Leutwyler, *Ann. Phys. (N.Y.)* **158** (1984) 142; *Nucl. Phys.* **B250** (1985) 465.
- [24] T. Becher and H. Leutwyler, Preprint BUTP-99-1.
- [25] M.A. Ivanov, J.G. Körner, V.E. Lyubovitskij, M.A. Pisarev, and A.G. Rusetsky, in preparation.
- [26] C. Caso et.al. (Particle Data Group), *Eur. Phys. J.* **C3** (1998) 1.

List of Tables

TABLE I Form Factors of $\Lambda_b \rightarrow \Lambda_c$ transitions.

TABLE II Form Factors of $\Lambda_c \rightarrow \Lambda_s$ transitions.

TABLE III Exclusive decay characteristics of Λ_b and Λ_c baryons.

TABLE I

Set	$q^2 = q_{max}^2$					
I	$F_1^V = 0.99$ (1)	$F_2^V = 0.034$ (0)	$F_3^V = 0.001$ (0)	$F_1^A = 0.97$ (1)	$F_2^A = -0.002$ (0)	$F_3^A = -0.035$ (0)
II	$G_1^V = 1.26$ (1)	$G_2^V = -0.20$ (0)	$G_3^V = -0.067$ (0)	$G_1^A = 0.96$ (1)	$G_2^A = -0.21$ (0)	$G_3^A = -0.076$ (0)
Set	$q^2 = 0$					
I	$F_1^V = 0.55$ (0.62)	$F_2^V = 0.017$ (0)	$F_3^V = 0.005$ (0)	$F_1^A = 0.54$ (0.62)	$F_2^A = -0.001$ (0)	$F_3^A = -0.017$ (0)
II	$G_1^V = 0.69$ (0.62)	$G_2^V = -0.10$ (0)	$G_3^V = -0.033$ (0)	$G_1^A = 0.54$ (0.62)	$G_2^A = -0.10$ (0)	$G_3^A = -0.036$ (0)

TABLE II

Set	$q^2 = q_{max}^2$					
I	$F_1^V = 0.70$ (0.62)	$F_2^V = -0.14$ (-0.044)	$F_3^V = -0.045$ (-0.044)	$F_1^A = 0.76$ (0.63)	$F_2^A = -0.032$ (-0.044)	$F_3^A = -0.16$ (-0.044)
II	$G_1^V = 1.16$ (0.77)	$G_2^V = -0.41$ (-0.20)	$G_3^V = -0.10$ (-0.20)	$G_1^A = 0.65$ (0.58)	$G_2^A = -0.45$ (-0.20)	$G_3^A = 0.15$ (-0.20)
Set	$q^2 = 0$					
I	$F_1^V = 0.41$ (0.33)	$F_2^V = -0.071$ (-0.018)	$F_3^V = -0.025$ (-0.018)	$F_1^A = 0.38$ (0.33)	$F_2^A = -0.021$ (-0.018)	$F_3^A = -0.084$ (-0.018)
II	$G_1^V = 0.65$ (0.39)	$G_2^V = -0.22$ (-0.08)	$G_3^V = -0.050$ (0)	$G_1^A = 0.34$ (0.31)	$G_2^A = -0.24$ (-0.08)	$G_3^A = 0.070$ (0)

TABLE III

Process	Heavy Quark Limit		Finite Quark Masses	
$\Lambda_b \rightarrow \Lambda_c + e^- \bar{\nu}_e$	$\Gamma = 5.4 \times 10^{10} \text{ sec}^{-1}$	$\alpha = -0.761$	$\Gamma = 4.9 \times 10^{10} \text{ sec}^{-1}$	$\alpha = -0.815$
$\Lambda_c \rightarrow \Lambda_s + e^+ \nu_e$	$\Gamma = 11.8 \times 10^{10} \text{ sec}^{-1}$	$\alpha = -0.798$	$\Gamma = 7.6 \times 10^{10} \text{ sec}^{-1}$	$\alpha = -0.878$

List of Figures

Fig. 1 Semileptonic decay of heavy baryon $\Lambda_{b(c)} \rightarrow \Lambda_{c(s)} + e + \nu_e$.

